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Cosmological Consequences of Nearly Conformal Dynamics at the TeV scale

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Abstract

Nearly conformal dynamics at the TeV scale as motivated by the hierarchy problem can be characterized by a stage of significant supercooling at the electroweak epoch. This has important cosmological consequences. In particular, a common assumption about the history of the universe is that the reheating temperature is high, at least high enough to assume that TeV-mass particles were once in thermal equilibrium. However, as we discuss in this paper, this assumption is not well justified in some models of strong dynamics at the TeV scale. We then need to reexamine how to achieve baryogenesis in these theories as well as reconsider how the dark matter abundance is inherited. We argue that baryonic and dark matter abundances can be explained naturally in these setups where reheating takes place by bubble collisions at the end of the strongly first-order phase transition characterizing conformal symmetry breaking, even if the reheating temperature is below the electroweak scale $\sim$ 100 GeV. We also discuss inflation as well as gravity wave smoking gun signatures of this class of models.

Contents

1 Introduction

Within the next few years, the LHC will be probing the electroweak (EW) symmetry breaking sector of the Standard Model (SM). In the SM, there is no understanding of the dynamics responsible for the breaking of the $SU(2)_L \times U(1)_Y$ gauge invariance. Electroweak symmetry breaking is introduced by hand through the addition of the Higgs mass operator.

There are two main avenues for explaining the lightness of the scalar Higgs boson: supersymmetry and Higgs compositeness. Minimal supersymmetry predicts a too light Higgs unless a tuning is invoked. The idea of Higgs compositeness has therefore received a revival of interest in the last few years [1,2]. In this framework, EW symmetry breaking is triggered by a spontaneously broken nearly scale invariant sector at a scale $\Lambda_{\mathrm{CFT}} \sim 4\pi f \geq \Lambda_{\mathrm{EW}} \sim 4\pi v$. The hierarchy $\Lambda_{\mathrm{EW}} \ll M_{\mathrm{Pl}}$ is explained dynamically via dimensional transmutation, as the quantum running of a dimensionless coupling generates a new scale, like in QCD. The presence of a moderate separation between the scale of EW symmetry breaking $v = 246$ GeV and the scale of conformal breaking f allows to keep under control the unwanted corrections to EW precision observables and is conceivable if the dynamics responsible for EW symmetry breaking is strongly coupled and nearly conformal, like in theories of walking technicolor [3] or via AdS/CFT in Randall-Sundrum extra-dimensional warped geometries [4–6]. The spectrum of states at the EW scale in these theories contains the pseudo-Goldstone dilaton from the spontaneous breaking of conformal invariance (the radion in the 5D picture).

Extensive studies have been devoted to the phenomenology of this scenario, in particular to make it consistent with experimental constraints and determine its distinctive signatures at colliders, see e.g. Ref. [7]. Somehow, the cosmological consequences of this framework have not been much explored although they can be completely different from the standard picture

[8–12] and open new and unique avenues for addressing the main cosmological puzzles such as the nature of dark matter, the origin of the matter-antimatter asymmetry and inflation[1]. The goal of this paper is to present some general properties which we believe are quite typical in a large class of theoretically well-motivated models. We will argue that in a certain class of theories of EW symmetry breaking with nearly conformal dynamics, we have the following:

1. A period of supercooling with some efolds of inflation due to a strongly first-order phase transition is likely.

2. The reheat temperature of the universe is then at the EW scale and depending on the mass of the radion and Higgs, it can even be below the typical sphaleron freeze-out temperature ($T_{\rm reh} \lesssim T_{\rm sph} \sim 120$ GeV.)

   Therefore, dilution of particle abundances during the short inflationary stage and a low reheat temperature both require to re-examine dark matter production and baryogenesis.

3. A large signal in gravity waves in the millihertz range is a smoking-gun signature of this scenario and therefore of nearly conformal dynamics at the TeV scale. For any other model with a strongly first-order phase transition at the EW scale, a detectable stochastic background of gravitational waves is fine-tuned.

The main distinctive feature of the cosmological scenario we are considering is two-fold: Reheating comes from bubble collisions following conformal symmetry breaking after the universe has undergone a stage of significant supercooling. Along, the reheat temperature of the universe is close to the scale of the induced EW symmetry breaking, and depending on its precise value, it will restrict the range of viable mechanisms for both visible and dark matter genesis. Our underlying framework is therefore very different from the common cosmological paradigm where the reheat temperature of the universe is assumed to be well above the EW scale, which on the other hand is rather natural if inflation occurs at a scale $\sim 10^{15}$ GeV. However, there is no real need for a reheat temperature much above the EW scale. As we will argue, it is possible to explain the visible and dark matter abundances even with a reheat temperature below the weak scale.

The paper is organized as follows. In Section 2, we present the main properties on which we rely and study the nature of the phase transition. We also estimate the number of efolds, summarize all the constraints and consider inflation. Section 3 is about reheating and consequences for baryogenesis and dark matter particle production. We discuss experimental probes in Section 4, and we conclude in Section 5.

2 A generic strongly first-order phase transition at the electroweak scale

In most of this paper, we will not make any particular assumptions about the nature of the strongly interacting sector. Our whole discussion will rely on the general assumption

[1]Literature exists on the possibility that dark matter is made of composite states or on baryogenesis proposals based on technibaryons, however it generally assumes a standard cosmological evolution, e.g. [13, 14].

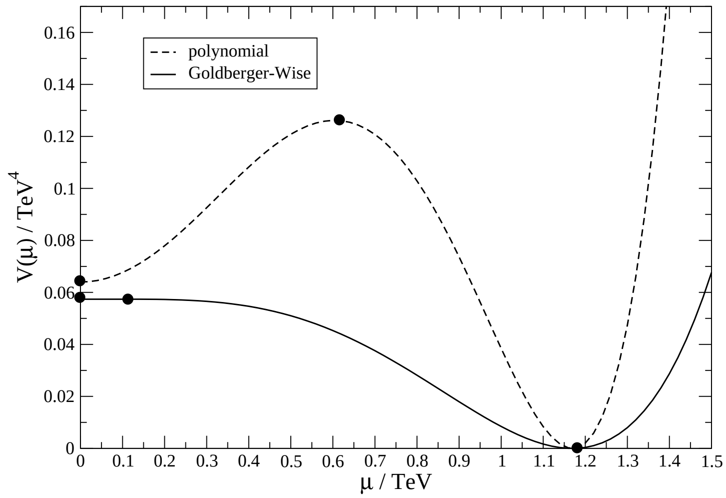

Figure 1: Comparison of a typical polynomial potential given here by $\lambda(\mu^2-\mu_0^2)^2+\frac{1}{\Lambda^2}(\mu^2-\mu_0^2)^3$ with a nearly conformal potential of the type of eq. (1). Both have a minimum at $\mu_{\rm min}\sim 1.2$ TeV. For the usual polynomial potential $\mu_{\rm max}/\mu_{\rm min}\sim\mathcal{O}(1)$, unless coefficients are fine-tuned while for the potential (1) with $|\epsilon|<1$, one can easily get a shallow potential with widely separated extrema. In this particular example $|\epsilon|=0.2$. The $\bullet$ indicates the position of the maxima.

that the scalar effective potential describing symmetry breaking is a scale invariant function modulated by a slow evolution:

$$V(\mu)=\mu^4 P\left[\,\left(\frac{\mu}{\mu_0}\right)^{\epsilon}\,\right], \tag{1}$$

similarly to the Coleman-Weinberg potential where a slow RG evolution of the potential parameters can generate very separated scales. P is a polynomial function reflecting some explicit breaking of conformal invariance by turning on some coupling of dimension $-\epsilon$. This potential generically has a minimum at $\mu_-\neq 0$. We are interested in the case where $|\epsilon|$ is small so that we have an almost marginal deformation of the CFT. If $\epsilon>0$ symmetry breaking results from a balance between two operators unlike in QCD where it is driven by the blow-up of the gauge coupling [5, 6]. For $|\epsilon|\ll 1$, a large hierarchy is generated.

2.1 Cosmological properties of a nearly conformal scalar potential

This class of potentials leads to some unique cosmological properties. In particular, it leads to a strongly first-order phase transition. What makes the nearly conformal potentials special is the fact that the positions of the maximum μ_+ and of the minimum μ_- can be very far apart in contrast with standard polynomial potentials where they are of the same order, as illustrated in Fig. 1. This makes the temperature dependence of the tunneling action behave very differently from the case of standard polynomial potentials. The nucleation temperature T_n is determined by the tunneling point μ_r (also called *release point*), which is located behind the barrier, somewhere between the maximum and the minimum of the potential. For a standard polynomial potential, μ_+ and μ_- are of the same order and the

tunneling point is of the same order as the value of the field at the minimum of the potential. For a nearly conformal potential, the two extrema are widely separated and as we will show, the release point can be as low as $\mu_r \gtrsim \sqrt{\mu_+\mu_-} \ll \mu_-$. Since the nucleation temperature $T_n \propto \mu_r$, we can get a very small T_n compared to the vacuum expectation value of the scalar field μ_- and therefore several efolds of inflation.

Typically, an extended phase of inflation (at least several efolds) cannot be ended by a first-order phase transition. This is the well-known graceful exit problem of old inflation which results from the following argument: for a generic free energy $V(\phi, T)$ the tunnel action S_3/T is a "well-behaved" (meaning roughly polynomial) function of the temperature T. The first nucleated bubbles appear when the temperature satisfies, in terms of the Hubble constant H,

$$S_3/T \approx \log \frac{T^4}{H^4}. \tag{2}$$

At the weak scale, this corresponds to $S_3/T \approx 140$. In order to realize several efolds of inflation, the onset of the phase transition and bubble nucleation should happen at a temperature T_n that is several orders of magnitude smaller than the critical temperature T_c defined as the temperature at which the symmetric and broken phase are degenerate.

If S_3 is a well-behaved function of T, characterized by the energy scale $\mu_0 \sim T_c$, its derivative $\partial_T(S_3/T)$ is likewise and the parameter β which quantifies the inverse duration of the phase transition satisfies

$$\beta/H = T \frac{d}{dT} \frac{S_3}{T}\bigg|_{T_n} \sim \frac{T_n}{\mu_0} \left.\frac{S_3}{T}\right|_{T_n}. \tag{3}$$

An extended phase of inflation (for example, $N_{\text{efolds}} \sim \log T_c/T_n \sim 10 \to T_n/T_c \sim 10^{-4}$) corresponds to $T_n \ll \mu_0$ then $\beta/H \ll 1$, which implies that bubbles never percolate and the phase transition cannot complete and reheating never occurs.

In contrast, the potential (1) leads to a tunneling action that is well-behaved as a function of μ^ϵ rather than μ. This way it is possible to achieve a small nucleation temperature together with bubble percolation and a rather long but finite duration of the phase transition for $\epsilon \sim \mathcal{O}(1/10)$

$$\beta/H = T \frac{d}{dT} \frac{S_3}{T}\bigg|_{T_n} \sim \epsilon \left.\frac{S_3}{T}\right|_{T_n} \gtrsim 1. \tag{4}$$

An example is given in Fig. 2 where the tunneling action is plotted for a specific Goldberger-Wise potential [15] (taken from Ref. [11]) in comparison with an action occurring e.g. in the electroweak phase transition in supersymmetric extensions of the SM.

Let us explain this more quantitatively. The conformal phase transition can be studied by working in a five-dimensional Anti de Sitter (AdS) space in which the radion is stabilized by a bulk scalar with a relatively small mass [8–11]. In the 4D picture, this corresponds to a balance between a marginal and a slightly irrelevant deformation of the gluon sector of the CFT. At high temperature, the system is in an AdS-Schwarzschild (AdS-S) phase involving a single ultraviolet (Planck) brane, providing the UV cutoff of the theory. The free energy of the AdS-S phase is given by

$$F_{\text{AdS-S}} = -4\pi^4 (Ml)^3 T^4, \tag{5}$$

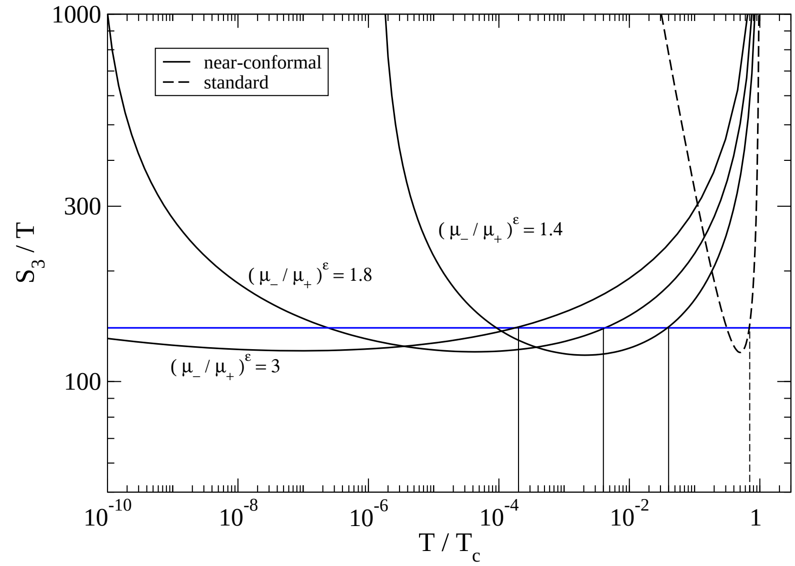

Figure 2: The tunneling action S_3/T as a function of T/T_c for a typical nearly conformal potential (solid line) (we used the Goldberger-Wise potential for illustration) and for a usual polynomial Higgs potential (dashed line). The horizontal blue line indicates the tunneling value $S_3/T \sim 4\log(M_{\rm Pl}/T_{\rm EW}) \sim 140$. For a standard potential, the nucleation temperature T_n is always close to the critical one, T_c, unless some fine-tuning is invoked. For a nearly conformal potential, supercooling is a general feature and T_n can easily be several orders of magnitude below T_c.

where l is the 5D AdS curvature, of the order of the Planck scale, and M the 5D Planck mass. By holography, $F_{\rm AdS-S}$ can be interpreted as the free energy of a strongly coupled large N CFT where the rank N of the $SU(N)$ dual gauge theory in four dimensions is related to $(Ml)^3$ via the AdS/CFT correspondence

$$N^2 = 16\pi^2 (Ml)^3. \tag{6}$$

Note that this relation (6) holds for the gauge/gravity duality in the $\mathcal{N} = 4$ supersymmetric context. In a more bottom-up approach to holography (like e.g. in AdS/QCD) the four dimensional theory has less supersymmetry and (6) is modified.

At low temperature, there are two branes, with a slice of AdS bulk in between. The infrared brane spontaneously breaks the conformal symmetry of the theory. The resulting effective potential of the radion is of the form (1) where the field μ is a reparametrization of the brane separation r

$$\mu = l^{-1} e^{-r/l}, \tag{7}$$

with a canonical kinetic term (up to a factor $12(Ml)^3$, see [11]). The position of the extrema $\mu_\pm$ of V depend on the specific parameters but are given by

$$\mu_+^\epsilon \lesssim \mu_-^\epsilon \lesssim 1. \tag{8}$$

In the Randall-Sundrum model, the smallness of $\epsilon \sim \mathcal{O}(1/10)$ is used to generate the hierarchy between the Planck and the electroweak scale, $\mu_- \ll l^{-1}$, but also implies $\mu_+ \ll \mu_-$ and the potential is nearly conformal between those widely spread values.

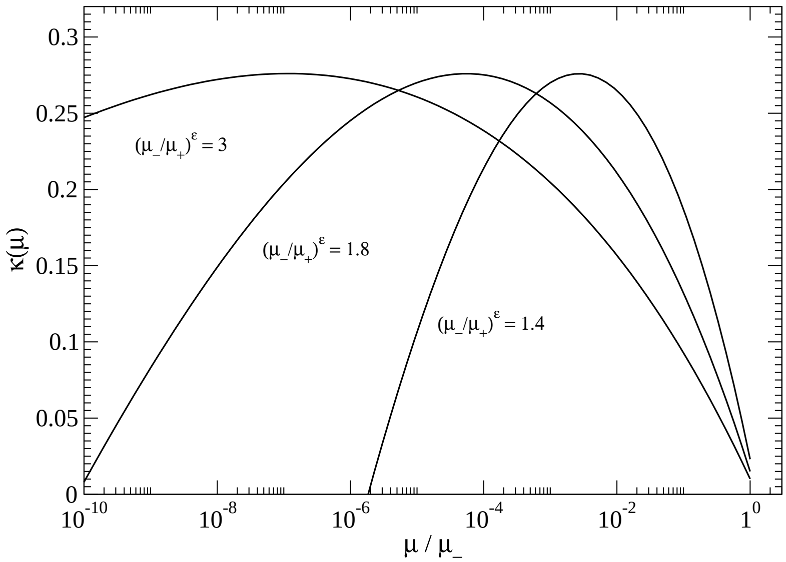

Figure 3: The running of the effective quartic coupling κ. Tunneling typically occurs when the quartic coupling is close to maximal, for a value of the release point close to $\mu_r \gtrsim \sqrt{\mu_+\mu_-}$, which can be orders of magnitude smaller than the value μ_- at the minimum of the potential.

The tunneling action can be calculated by determining the bounce solution for the potential (1) [8,9]. An accurate approximation can be obtained by exploiting the nearly conformal behavior of the system (we follow the notation and analysis of [11]). For a certain bounce solution with release point μ_r, the potential is approximated by

$$V(\mu) \approx \mu^4 P((\mu_r/\mu_0)^\epsilon) \equiv -\mu^4 \kappa. \tag{9}$$

The conformal invariance of the potential then allows to determine the action and the corresponding nucleation temperature T_n as (we only consider the $O(3)$ symmetric tunnel action here)

$$S_3/T \simeq 290\kappa^{-3/4}(Ml)^3, \quad T_n \simeq 0.1\kappa^{1/4}\mu_r. \tag{10}$$

The action of a critical bubble is minimal when the quartic coupling κ is maximal. Quite generally, the value of κ is bounded by a value around $\frac{1}{2}$ [11]. Hence, for large values of $(Ml)^3$, the phase transition is strong and the symmetric phase can even become stable [8]. Accordingly, for the phase transition to complete, we must have

$$(Ml)^3 \lesssim 0.3\,. \tag{11}$$

We will display more precise conditions on the parameter space in the next section.

The function S_3/T is shown in Fig. 2 and $\kappa(\mu)$ is plotted in Fig. 3. For most of the parameter space, the system tends to tunnel close to the maximum of κ. Since κ is a second order polynomial function of $\log\mu$ in the limit $\epsilon \ll 1$, its maximum is given by $\log\mu_r \gtrsim (\log\mu_+ + \log\mu_-)/2$, hence

$$\mu_r \gtrsim \sqrt{\mu_-\mu_+} \ll \mu_-, \quad T_n \ll T_c. \tag{12}$$

Therefore, the Universe tends to be cold at the onset of the phase transition. On the other hand, for the duration of the phase transition one finds

$$\beta/H \;=\; T\frac{d}{dT}\frac{S_3}{T}\bigg|_{T_n} \approx \frac{3}{4}\frac{S_3}{T}\left(\mu_r\frac{d}{d\mu_r}\kappa\right) \sim \frac{3\epsilon}{4}\frac{S_3}{T}, \tag{13}$$

and percolation is not an issue for $\epsilon \gtrsim 10^{-2}$.

Next, we investigate what are the typical expectations for the number of efolds and what controls the amount of supercooling. For that, we use the explicit Goldberger-Wise potential.

2.2 Typical amount of supercooling (number of efolds)

The number of efolds of inflation is given by

$$N_{\text{efolds}} \sim \log\frac{T_c}{T_n}. \tag{14}$$

In this section we estimate the two temperatures T_c and T_n and the typical amount of supercooling of the conformal phase transition. As we have seen previously, the nucleation temperature T_n is proportional to the release point μ_r :

$$T_n \simeq 0.1\,\mu_r\,\kappa(\mu_r)^{1/4}, \tag{15}$$

which is bound by the position of the maximum of the quartic coupling $\kappa(\mu)$ given by the scale $\sqrt{\mu_-\mu_+}$. In the Goldberger-Wise mechanism of the Randall-Sundrum model, parameters of the potential are fixed such that the ratio μ_-/μ_+ explains the hierarchy between the EW scale and the Planck scale. Following the notations used in [11], the ratio μ_-/μ_+ can be written as

$$\frac{\mu_-}{\mu_+} = \left(\frac{\xi_-}{\xi_+}\right)^{1/\epsilon} \quad \text{where} \quad \xi = \frac{v_1}{v_2}e^{-\epsilon r/l}, \tag{16}$$

and v_1 and v_2 are the expectation values of the Goldberger-Wise bulk scalar field at the two boundaries. The parameters ξ_-, ξ_+ and ϵ are related by (see eq. (49) of [11])

$$\xi_+ \;=\; \frac{4+\epsilon}{2+\epsilon} - \xi_- \approx 2 - \xi_- \quad \text{for small detuning on the brane} \tag{17}$$

$$\xi_+ \;\approx\; \xi_-^2\frac{v_2}{v_1} \quad \text{for large detuning on the brane} \tag{18}$$

The detuning is to a modification of the values of the two brane tensions that are adjusted to ensure a vanishing 4D cosmological constant in the absence of back reactions. In particular one always has

$$\xi_+ > \xi_-^2\frac{v_2}{v_1}. \tag{19}$$

We therefore obtain from the EW/Planck scale hierarchy the upper bound

$$\frac{\mu_-}{\mu_+} = \left(\frac{\xi_-}{\xi_+}\right)^{1/\epsilon} < e^{r_-/l} \approx e^{37} \approx 10^{16}, \tag{20}$$

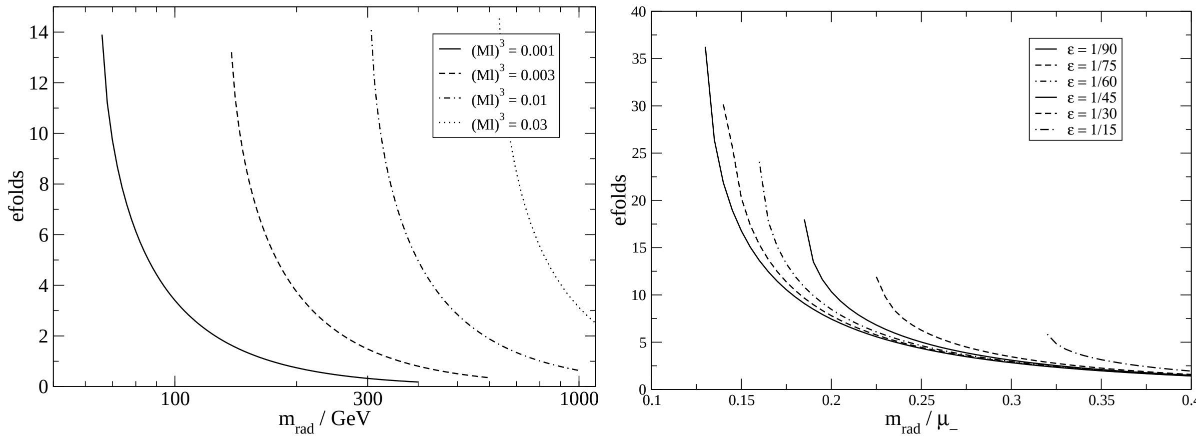

Figure 4: Number of efolds of inflation as a function of the radion mass. Left: Randall-Sundrum model for different values of $(Ml)^3$ and $\mu_- = 4$ TeV; Right: A generic model with potential (1) where the constraint fixing the hierarchy, eq. (20), is relaxed. At the point where the curves stop, the system cannot tunnel and is stuck in the symmetric phase.

which in turn gives a lower bound on the nucleation temperature according to (12).

As for the critical temperature T_c, it is given by equating the free energy of the AdS-S phase, with the potential difference between the conformally symmetric and broken phases. This potential difference is given by

$$\Delta V = 2\epsilon l^{-4}(Ml)^3 \frac{v_2^2}{M^3}\xi_-(\xi_- - 1)e^{-4r_-/l}, \tag{21}$$

and using the expression for the radion mass (eq. (69) of [11])

$$m_{\rm rad}^2 = \frac{2}{3}\epsilon\mu_-^2 \frac{v_2^2}{M^3}(4(\xi_- - 1) + \epsilon\xi_-)\xi_-, \tag{22}$$

we have for ξ_- not too close to unity

$$\Delta V \approx \frac{3}{4}(Ml)^3 m_{\rm rad}^2 \mu_-^2. \tag{23}$$

Equating this with the free energy of the AdS-S phase, eq. (5), yields

$$4\pi^4 T_c^4 = \frac{3}{4} m_{\rm rad}^2 \mu_-^2, \tag{24}$$

leading to

$$N_{\rm efolds} \sim \log\frac{T_c}{T_n} \simeq \log\frac{\mu_-}{\mu_r} + 0.74 - \frac{1}{4}\log\kappa(\mu_r) + \frac{1}{2}\log m_{\rm rad}/\mu_-. \tag{25}$$

Therefore, the number of efolds is essentially controlled by the hierarchy μ_-/μ_+ and the maximal number of efolds is

$$N_{\rm efolds} < \log\sqrt{\frac{\mu_-}{\mu_+}} = \frac{r_-}{2l} \approx 18. \tag{26}$$

We plot in Fig. 4(a) the number of efolds (more precisely $\log \mu_-/\mu_r$) as a function of the radion mass in the Randall-Sundrum model where the ratio μ_-/μ_+ is constrained by the weak/Planck scale hierarchy. This constraint is relaxed in Fig. 4(b) where therefore the number of efolds can reach values required to solve the horizon problem. In any case, we see that a large number of efolds is associated with a light radion (relative to the scale μ_-).

2.3 Backreaction constraints

In this section, we derive the limits to the validity of our analysis, in particular the constraints from backreaction. In the Goldberger-Wise stabilization mechanism [15] that leads to a potential of the form (1), one introduces a 5D bulk scalar Φ with a mass m that is related to ϵ by

$$\epsilon = \sqrt{4 + m^2 l^2} - 2. \tag{27}$$

As discussed in [9], imposing that the energy in the Goldberger-Wise field is subdominant compared to the bulk cosmological constant (in other words that it does not distort too much the AdS geometry) significantly restricts the available parameter space.

The 5D metric is parametrized as

$$ds^2 = e^{2A(r)} \left(dt^2 - e^{2\sqrt{\Lambda}t} d\vec{x}^2 \right) - dr^2, \tag{28}$$

where Λ is the 4D cosmological constant. Our parameter space comprises the radion mass m_{rad} , ϵ and $(Ml)^3$ (or N in the 4D language) and is constrained by requiring that some terms in the equation of motion of the warp factor A are small:

$$A'^2 = \frac{1}{l^2} - \frac{1}{24M^3} m^2 \Phi^2 + \frac{1}{24M^3} \Phi'^2 + \Lambda e^{-2A}. \tag{29}$$

Demanding that the first term in (29) dominates over the second and third leads to

$$4\epsilon v_1^2 \ll 24M^3, \quad \text{and} \quad v_2^2 (\epsilon \xi_- + 4(\xi_- - 1))^2 \ll 24M^3. \tag{30}$$

In addition, we should consider the impact of the 4D cosmological constant term Λ on the parameter space. If it is neglected, the parameters ξ_- and ξ_+ are related by (17) while in the regime of a large cosmological constant one finds the relation (18) and in general the bound (19). Interestingly, this bound automatically ensures percolation, c.f. eq. (4). Using the radion mass (22) and the relation $\xi_- = \frac{v_1}{v_2} e^{-\epsilon r_-/l}$, the three constraints are

$$\begin{aligned} m_{\rm rad}^2/\mu_-^2 &< 16\epsilon \frac{\xi_-}{4(\xi_- - 1) + \epsilon \xi_-}, \\ m_{\rm rad}^2/\mu_-^2 &< 4e^{-2\epsilon r_-/l} \frac{4(\xi_- - 1) + \epsilon \xi_-}{\xi_-}, \\ \epsilon &> \frac{l}{r_-} \log \xi_-/\xi_+. \end{aligned} \tag{31}$$

These constraints together with the contour lines for the predicted number of efolds are shown in Fig. 5 for three values of $(Ml)^3 = N^2/16\pi^2$. These plots clearly demonstrate the

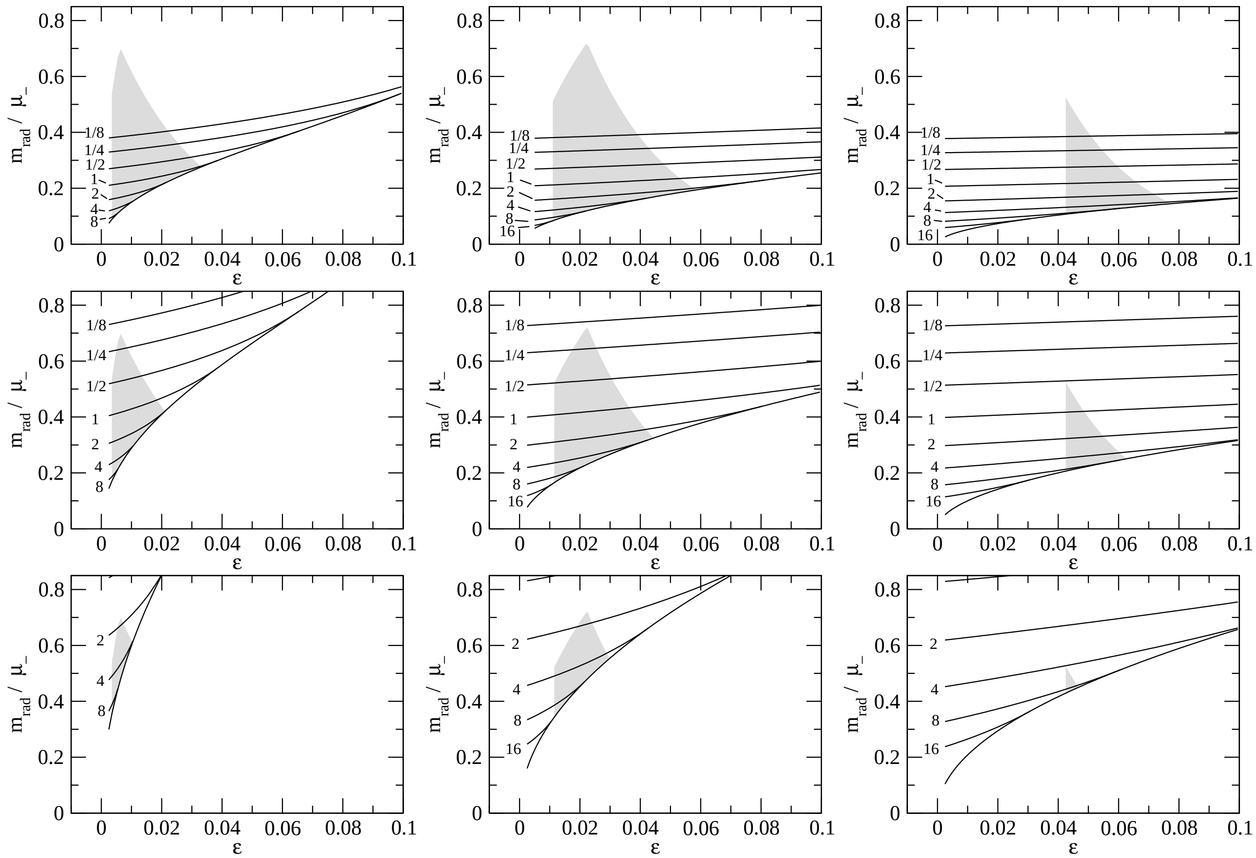

Figure 5: Contours for the number of efolds (more precisely $\log \mu_-/\mu_r$). The shaded region is where calculability can be trusted, as defined by the constraints in (31). Below the bottom line, the system never tunnels to the broken phase. From top to bottom, the plots show $N = 2$, 3 and 5; from left to right the series of three plots respectively use $\xi_-/\xi_+ = 1.05$, 1.2 and 1.65. For larger N the phase transition becomes stronger and beyond $N > 6$ the system is generally stuck in the symmetric phase, at least in the domain of calculability and only considering thermal tunneling.

observation first made in Ref. [8] that there is a tension between the large N assumption needed for calculability and the possibility to complete the phase transition. In [8], the bounds were stronger and the conclusion was rather negative, i.e that the transition could not complete in the regime of calculability. This conclusion was ameliorated in [9] where the tunneling action was estimated in the supercooling regime, namely in the thick-wall limit, and for $O(4)$ symmetric bubbles and also taking into account the fact that the field value at tunneling is not close to the value at the minimum. These effects improve the nucleation probability, as re-examined in more details in [10], and refined by taking into account backreactions in [11], confirming that the bounds are actually less stringent than in [8]. One also gets a much weaker phase transition when the geometry is deformed in the infrared [16].

To conclude, in the region of parameter space allowing calculability, the phase transition tends to be so strong that several efolds of inflation can occur before the onset of the phase transition. If one is willing to push calculability to its limit, $N = 2$, then there is typically less supercooling and several efolds take place only by tuning the radion mass to a low value.

Finally, note that in this paper, we present results only for $\epsilon > 0$. It would be interesting to consider in more details the $\epsilon < 0$ case which was argued to be more promising in [9], although in this case strong coupling effects arise in the IR and may be important before nucleation, thus preventing a reliable analysis. Furthermore, our discussion has been based on the Goldberger-Wise potential. There are alternative stabilization mechanisms relying on the Casimir forces induced by bulk fields [17]. Although we do not expect that the picture we have exposed would be qualitatively different, it would certainly be interesting to study the phase transition in this context.

2.4 Graceful first-order inflation from nearly conformal dynamics

Guth's original idea of inflation [18] was precisely that inflation could end by the tunneling of the false vacuum into the true vacuum during a first-order phase transition. However, it was realized that true vacuum bubbles would not percolate to give rise to the primordial plasma of relativistic degrees of freedom [19]. This drawback was solved in slow-roll models of inflation [20, 21] where a scalar field, the inflaton, slowly rolls down along its very flat potential.

As shown in the previous section, with a nearly conformal potential (1) it is possible to have a stage of inflation ended by a first-order phase transition, with no particular fine-tuning in the potential. However, it is clear from fig 4(a) that we cannot reach naturally a sufficiently large number of efolds to solve the standard cosmological problems that inflation is supposed to solve, even though for an inflationary stage taking place at the electroweak scale, we need only ~ 30 efolds to solve the horizon problem rather than ~ 60 if the inflation scale M_I is at the GUT scale, according to

$$N_{\rm efolds} = 62 - \log\left(\frac{10^{16}\,\mathrm{GeV}}{M_I}\right) - \frac{1}{3}\log\left(\frac{M_I}{T_{\rm reh}}\right). \tag{32}$$

Besides, even if we tune the radion mass so that a large number of efolds is achieved, one still has to solve the problem of generation of density perturbations as reheating from bubble collisions can only produce isocurvature density perturbations [22] and an extra source is needed to explain the generation of density perturbations. Anyhow, we note that there are no constraints on our extended stage of supercooling from either Big Bang Nucleosynthesis or from the Cosmic Microwave Background [22]. Imposing that there is no large fraction of bubbles that are Hubble size when nucleosynthesis commences and similarly that there is no large inhomogeneous regions near the last scattering surface that would lead to distortions in the CMB, results in weaker conditions than our criterion for percolation (4).

3 Reheating temperature and implications for baryogenesis and dark matter

Phenomenological consequences of a strongly first-order phase transition have been extensively discussed in the literature. We stress here that we distinguish the phase transition associated with conformal symmetry breaking from the one associated purely with EW symmetry breaking, although they are intertwined. The full EW symmetry breaking sector has

a potential of the form

$$V_{\rm TOT} = \mu^4 \times \left(P((\mu/\mu_0)^\epsilon) + \mathcal{V}(\phi)/\mu_0^4 \right). \tag{33}$$

While μ condensation induces EW symmetry breaking (and thus bubbles also involve Higgs field varying vev), one should keep in mind that the potential $\mathcal{V}(\phi)$ alone may not necessarily display a first-order phase transition. On the other hand, it is the phase transition associated with ϕ condensation which is traditionally the relevant one for baryogenesis. The nature of the phase transition in composite Higgs models remains to be investigated in specific models. Studies of the scalar potential have concentrated on $\mathcal{V}(\phi)$ [1, 23]. However, one should in principle compute the full $V_{\rm TOT}(\mu, \phi)$, which is a non-trivial task. Although this is a model-dependent question that relies on the form of the Higgs potential $\mathcal{V}(\phi)$ which we do not specify here, some general statements can be made on the cosmology as we discuss now.

3.1 Reheating temperature predictions

In our scenario, when bubbles are nucleated, the universe is very cold. Reheating starts when bubble collide. The value of the reheat temperature is a crucial ingredient to determine not only what are the possible frameworks for baryogenesis but also what are the underlying conditions for the computation of the dark matter abundance. Baryogenesis depends on whether the reheat temperature $T_{\rm reh}$ is below or above the sphaleron freeze-out temperature, which is essentially given by the temperature at which the electroweak symmetry is broken. As far as dark matter is concerned, if $T_{\rm reh}$ is at the electroweak scale, the common thermal freeze-out mechanism prediction for WIMP dark matter abundance is no more guaranteed if $T_{\rm reh} \lesssim m_{\rm DM}$.

The process of reheating from bubble collisions was first discussed in [24] and later in the early nineties in [25–27]. We apply these results to the case of a strongly first-order phase transition taking place in an empty universe as a result of nearly conformal dynamics at the TeV scale in the companion paper [28], where, in particular, we argue that the scalar field dynamics during the collisions provides ideal conditions for a natural *cold* baryogenesis mechanism. In the present section, we discuss the value of the reheat temperature $T_{\rm reh}$ and review the consequences for baryogenesis and dark matter generation mechanisms in this new context.

At the TeV scale, the expansion of the Universe is negligible and $T_{\rm reh}$ can be estimated using energy conservation

$$\Delta V = g^* \frac{\pi^2}{30} T_{\rm reh}^4, \tag{34}$$

where $g^* \sim 100$ is the number of relativistic degrees of freedom given by the particle content of the SM after the phase transition. On the other hand, the critical temperature is [8–11]

$$4\pi^4 (Ml)^3 T_c^4 = \Delta V. \tag{35}$$

Note that in contrast with more common phase transitions, the reheating temperature may therefore exceed the critical temperature associated with conformal symmetry breaking, if there is a large number of degrees of freedom in the CFT gas:

$$(Ml)^3 \gtrsim \frac{g_*}{120\pi^2} \sim 8 \times 10^{-2} \quad \text{or} \quad N > \sqrt{\frac{2g_*}{15}} \sim 3.6\,. \tag{36}$$

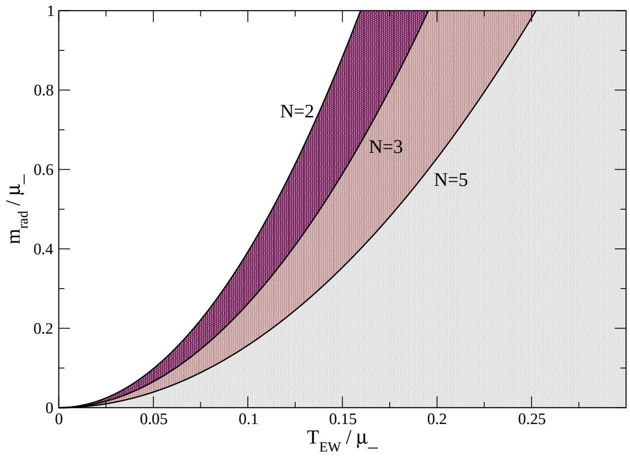

Figure 6: Value of the radion mass (in the colored regions) for which the reheat temperature is below the temperature at which the EW symmetry is restored, $T_{\rm EW}$, according to eq. (39).

The condition (36) is only indicative. First of all, after conformal symmetry breaking, most degrees of freedom of the conformal sector have masses comparable to the temperature and will, to a certain extent, contribute to the free energy, modifying the relation (34). Secondly, for a nearly conformal radion potential, the tunneling back to the symmetric phase involves sizable superheating and happens at a significantly larger temperature than the critical one (just as the tunneling to the broken phase involves sizable supercooling). In any case, even if nucleation back to the symmetric phase is a priori possible, bubbles of symmetric phase cannot grow rapidly into the broken phase, as the latent heat is negative (see e.g. [29]) and the phase transition has to proceed by other means (e.g. a slow growth of droplets). In summary, we do not need to consider any particular constraint resulting from a reheating temperature potentially higher than the critical temperature associated with conformal symmetry breaking.

What we are really interested in is whether the electroweak symmetry can be restored even if the conformal symmetry stays broken after reheating. Although this essentially depends on the form of the effective potential for the Higgs field that we have not specified here, some general conclusions can be drawn. For instance, if the reheat temperature is smaller than $T_{\rm EW}$, the temperature at which the EW symmetry is restored, then standard EW baryogenesis cannot take place and we have to rely on a different mechanism to explain the matter-antimatter asymmetry of the universe. Alternative mechanisms exist and will depend on the reheating temperature that we estimate now. Using again the following expression for the free energy difference in terms of the radion mass

$$\Delta V = \frac{3}{4}(Ml)^3 m_{\rm rad}^2 \mu_-^2, \tag{37}$$

and (34) we obtain for the reheating temperature

$$T_{\rm reh} = \left(\frac{45}{2\pi^2 g_*}\right)^{1/4} (Ml)^{3/4} \sqrt{m_{\rm rad}\mu_-}, \tag{38}$$

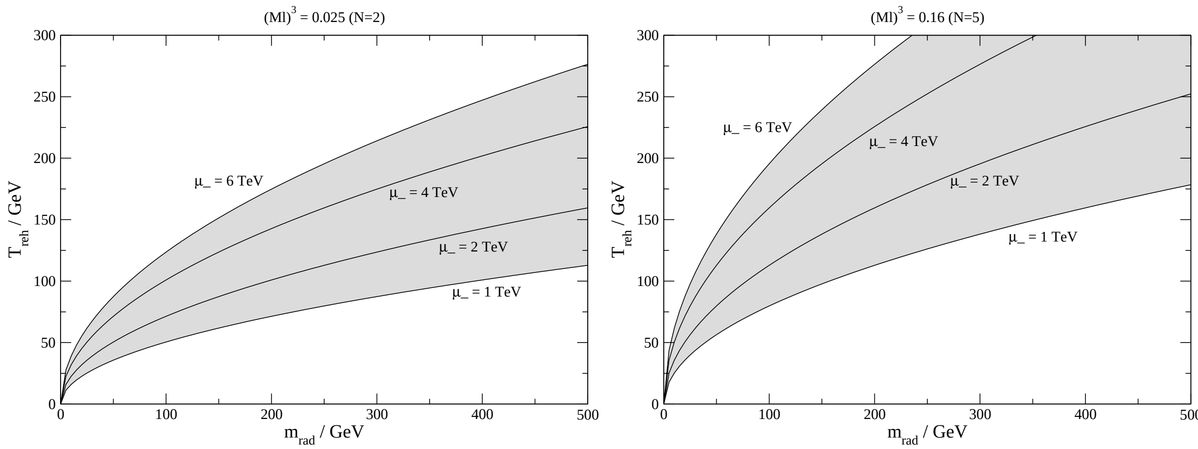

Figure 7: Reheat temperature $T_{\rm reh}$ as a function of the radion mass for different values of the scale of conformal symmetry breaking μ_- and two values of $(Ml)^3$, according to eq. (39).

which leads to

$$T_{\rm reh} \gtrless T_{\rm EW} \quad \longrightarrow \quad \frac{m_{\rm rad}}{\mu_-} \gtrless \frac{6.6}{(Ml)^{3/2}} \left(\frac{T_{\rm EW}}{\mu_-}\right)^2 . \tag{39}$$

The bounds on $m_{\rm rad}/\mu_-$ in the Goldberger-Wise model are shown in Fig. 5 and the bound (39) is illustrated in Fig. 6. We plot the reheat temperature as a function of the radion mass in Fig. 7. Thus, whether the electroweak symmetry is restored after reheating depends on the radion mass and on the temperature of electroweak symmetry breaking which in turn depends on the Higgs mass.

For example, in the minimal composite Higgs model [1], rather large Higgs masses can arise, in particular at small N, while for the temperature of electroweak symmetry breaking, the Standard Model relation [30] is valid

$$\frac{T_{\rm EW}}{M_{\rm Higgs}} \simeq \left(m_W^2/v^2 + \frac{1}{2} m_Z^2/v^2 + m_t^2/v^2 \right)^{-1/2} \sim 1.3 \, . \tag{40}$$

In the case $N = 3$, significant supercooling happens if $m_{\rm rad}/\mu_- < 0.3$ and for a Higgs mass of 200 GeV, the electroweak symmetry is restored after reheating if $\mu_- > 2.6$ TeV. Similarly, one finds in the case $N = 5$ the bound $\mu_- > 1.6$ TeV.

3.2 Viable baryogenesis mechanisms

• $T_{reh} > T_{EW}$: If the reheat temperature is large enough, the universe will go back temporarily into the electroweak symmetric phase. In this situation, we will recover a standard cosmological evolution where eventually the EW symmetry is broken.

If the EW symmetry breaking proceeds by a cross-over, no departure from equilibrium occurs and none of the common baryogenesis mechanisms can apply here (high scale leptogenesis or EW baryogenesis). Alternatives could be baryogenesis from out-of-equilibrium decay of TeV scale composite states [31] or low-scale leptogenesis with TeV scale particles that produce the lepton asymmetry by decay. These TeV scale particles can be produced by

bubble collisions [27, 28]. In the case of a Yukawa coupling between the fermionic species ψ and the scalar field ϕ

$$\mathcal{L} = y\,\phi\,\bar{\psi}\,\psi, \tag{41}$$

a fraction y^2 of the energy of the scalar sector is released into the fermions. In particular, for relativistic bubble wall velocities, even particles with masses far beyond the electroweak scale can be produced in bubble collisions [27, 32, 33]. In our scenario, almost all the energy resides in the scalar sector at the time of bubble collisions [28], therefore, only very moderate couplings can produce sizable particle numbers. It is thus tempting to consider the possibility of TeV scale leptogenesis. Nevertheless, a sufficient production of right-handed neutrinos requires Yukawa couplings that are in conflict with the generation of light neutrino masses by the seesaw mechanism

$$y^2 \sim \frac{m_\nu m_N}{v^2}. \tag{42}$$

In this case, the number of right-handed neutrinos can be estimated (ρ and s denote the total energy and entropy densities)

$$n_N \sim y^2 \frac{\rho}{m_N} \sim \rho\,\frac{m_\nu}{v^2} \qquad \to \qquad \frac{n_N}{s} \ll 10^{-10}. \tag{43}$$

Such value prevents a successful leptogenesis even assuming maximal CP-violating effects. Thus, the lepton asymmetry has to be produced from e.g. the decay of a new TeV species that is not responsible for the light neutrino masses [34]. In our scenario, this is especially easy because new TeV particles from the strongly interacting sector with $\mathcal{O}(1)$ couplings are abundantly produced non-thermally and their subsequent decays happen far out-of-equilibrium such that washout is avoided.

If the EW phase transition is first-order, standard EW baryogenesis is in principle also possible and it will be interesting to investigate further whether composite Higgs models offer the necessary conditions (for an effective description see [35, 36]). In particular, the extended Higgs sector in [23] seems promising. Note however that for too strong EW phase transitions, supersonic bubbles suppress CP violating densities in front of bubble walls, thus preventing the mechanism of EW baryogenesis to work.

• $\underline{T_{reh} < T_{EW}}$: If on the other hand the reheat temperature is below the temperature at which the EW symmetry is restored, a very interesting and non-trivial possibility naturally opens up: *cold* electroweak baryogenesis, which is the subject of our companion paper [28].

We summarize the various baryogenesis possibilities in Table 1.

3.3 Dark matter production during reheating by bubble collisions

The present scenario of conformal and electroweak symmetry breaking has also important implications for the dark matter abundance. A stage of inflation at the onset of the phase transition dilutes not only any preexisting baryon asymmetry but also the abundance of dark matter. As long as this era of inflation only lasts a couple of efolds, this dilution effect is not too severe and standard scenarios of dark matter, namely the WIMP scenario, could account for the observed dark matter. However, beyond around eight efolds, the preexisting dark matter must have been overabundant before the electroweak phase transition in order to be in accordance with today's observations, which is rather implausible.

	$T_{\rm reh} > T_{\rm EW}$		$T_{\rm reh} < T_{\rm EW}$	
	EWPT is 1st-order	EWPT is crossover	$\frac{\phi}{T}\Big\|_{T_{\rm reh}} > 1$	$\frac{\phi}{T}\Big\|_{T_{\rm reh}} < 1$
cold EW baryogenesis	$-$	$-$	$+$	$-$
non-local EW baryogenesis	if $\phi/T\|_{\rm EW} > 1$	$-$	$-$	$-$
low-scale lepto/baryogenesis from TeV particle decays	$+$	$+$	$-$	$+$
B-conserving baryogenesis from asymmetric dark matter	$+$	$+$	$+$	$+$

Table 1: The viability of different baryogenesis scenarios depending on the reheating temperature after the conformal phase transition and the properties of the electroweak phase transition. T_{EW} is the temperature at which EW symmetry gets restored. The last possibility is very specific as it does not require sphaleron processes at any moment. It assumes that, in a B-conserving universe, the dark matter carries the anti-baryonic charge that is missing in the visible sector [31].

There is actually no dilemma. As discussed above and in [28], any TeV-mass particle with significant coupling to the radion/Higgs, hence possibly dark matter particles, will be substantially produced at the early stages of preheating when bubbles collide. With a Yukawa coupling y, a fraction y^2 of the energy in the scalar sector is transformed into dark matter particles and already very moderate couplings of order 10^{-5} between the radion/Higgs and the dark matter sector (that contains for example stable composite states of the strongly coupled sector) will account for the observed dark matter abundance [32, 33].

4 Experimental probes

4.1 LHC tests

The underlying motivation for the framework we have been discussing is, as well-known, that strong dynamics at the TeV scale nullifies the hierarchy problem. The standard realization of this scenario is technicolor [37], which, however, is not easy to reconcile with EW precision measurements and flavor constraints. In the last years, an interesting variation interpolating between technicolor theories and the SM Higgs model has appeared where the Higgs emerges as a pseudo Goldstone boson from the breaking of a global symmetry of a strongly interacting sector [2]. Generically, in this picture, we expect new resonances at the scale $\mu_- \sim \mathcal{O}(1)$ TeV. However, depending on the precise value μ_-, these states may or may not be accessible at the LHC. A genuine strong coupling signature is the growth with energy of the longitudinal gauge bosons scattering amplitudes and double Higgs production. Observing these effects has been shown to be extremely challenging and would require several hundreds of fb^{-1} of LHC data at 14 TeV [38].

On the other hand, these models also suggest that SM fermion masses should arise via mixing of elementary fermions with composite fermions of the strong sector [7]. In this

context, the top quark is mainly a composite object while the other light SM quarks are mainly elementary. A natural prediction is then the existence of light ($\lesssim 1$ TeV) fermionic composite partners of the third generation fermions, in particular the top quark [39–41]. At the LHC, composite quarks can be pair-produced with a large QCD cross section. They can also be singly produced. Prospects at the LHC for their discovery are very promising [42,43]. There is a large number of phenomenological studies related to this class of models, for instance related to four-top events [44, 45] or signatures associated with composite leptons [46].

Besides, the rate for Higgs production and decay can significantly differ from the SM prediction. Depending on the choice of parameters, in particular on the ratio v/f, Higgs searches may either be deteriorated or the Higgs production rate may be enhanced with respect to the SM [47]. Moreover, non-minimal composite Higgs realizations typically lead to a multi-Higgs framework. The complexity of the scalar sector depends on the global symmetry of the strong sector [23].

Finally, we note that a light dilaton, as the pseudo-Goldstone boson of spontaneously broken scale invariance, can fake the Higgs at the LHC. Distinguishing it from a minimal Higgs is a subtle issue [48, 49].

4.2 A smoking gun stochastic gravity wave signal

Gravity wave signals from the conformal phase transition have been studied in [9, 11]. The stronger is the transition, the larger is the latent heat release and therefore the larger is the amplitude of the gravity wave signal. The size of the signal generally scales as $(\beta/H)^{-2}$. It should then be clear from our discussion in Section 2 that an observable signal cannot be obtained without fine-tuning in the case of an ordinary polynomial phase transition, as well illustrated by Fig. 5 of [36]. However, for a nearly conformal Goldberger-Wise framework, several factors favor the production of a stochastic gravity wave background that could be observed with space-based interferometers such as LISA:

- Due to the large supercooling, almost all of the energy of the system resides in the bubble walls during the phase transition. This kinetic energy is essential to give rise to a sizable anisotropic stress that produces gravitational radiation [9].

- The phase transition proceeds rather slowly, c.f. eq. (4), such that typically $\beta/H \gtrsim O(10)$ while for generic strong phase transitions one finds $\beta/H \gtrsim O(100)$.

- The peak frequency of the gravity wave spectrum is proportional to the temperature of the phase transition. Since in the present scenario the phase transition results from the breaking of the conformal symmetry and hence dynamics at the TeV scale, the peak frequency is by about one order of magnitude larger than for a generic electroweak phase transition [50]. This improves the prospects for observation at LISA.

In conclusion, the observation of a gravity wave spectrum peaked in the millihertz range would indicate either some sort of conformal dynamics at the TeV scale, from a strongly interacting sector as in the Randall-Sundrum setup (see also [51]), or some form of low scale inflation [52].

Finally, note that while this paper has focused on the TeV scale, the discussion can be applied to any other scale as the properties of the phase transition do not depend on the absolute energy scale but only on the amount of supercooling. For instance, nearly conformal dynamics at an intermediate scale ($\sim 10^7$ GeV), would lead to a gravity wave spectrum peaked in the 10-100 Hz range, and could thus be probed by LIGO [53].

5 Conclusion

The framework in which EW symmetry breaking is triggered by a strongly coupled nearly conformal sector offers an appealing dynamical solution to the hierarchy problem. It will take some time at the LHC to determine whether the origin of the EW scale is due to a new strong sector. Somehow, the cosmological consequences associated with this scenario have not been much explored. In this paper, we have stressed some peculiar properties of the phase transition associated with conformal symmetry breaking and provided a study of possible interesting cosmological features by making the least possible reference to explicit models.

We have shown how a nearly conformal potential can lead naturally to a significant period of supercooling. Any ordinary polynomial potential has to be fine-tuned to lead to several efolds of inflation ended by a first-order phase transition or the latter never completes. With a potential (1), there is no eternal inflation problem as bubbles can percolate and reheat the universe. Although the number of efolds is moderate and not sufficient to solve the horizon problem, there are still important consequences of phenomenological interest. While we have used the Goldberger-Wise potential for illustration, the qualitative features we have outlined are general, in particular:

- A strongly first-order phase transition
- Reheating from bubble collisions
- A reheat temperature possibly below the sphaleron freeze-out temperature
- Efficient out-of-equilibrium heavy particle (or classical field configuration) production
- A smoking gun gravity wave stochastic background peaked in the millihertz range

Heavy particle production from bubble collisions was already studied in details in the nineties. We find it somewhat appealing that the framework we are proposing here makes this possibility quite natural and motivates alternative cosmological scenarios from the standard one. We refer the reader to [28] for a more detailed discussion on reheating during bubble collisions at the TeV scale where in particular we advocate a large production of Higgs winding configurations and provide a description of the cold baryogenesis mechanism in this context.

Acknowledgments

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